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# LATEST DEVELOPMENTS ON THE BESS MODEL

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**Abstract**

The idea of a strongly interacting sector as responsible for the electroweak symmetry breaking is tested through an effective lagrangian description, called the BESS model, constructed on the standing point of custodial symmetry and gauge invariance, without specifing any dynamical scheme.

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# 1 Introduction

Although the standard model of electroweak interactions (hereafter denoted as SM) is in perfect agreement with experimental data — at a precision level of below the 0.5% — the mechanism of electroweak symmetry breaking (hereafter denoted as EWSB) is still an open question and certainly the central issue and the main goal for future accelerators.

In our opinion one possibility is a perturbative scheme with a light Higgs which has still to be discovered. It needs a new symmetry like supersymmetry to be viable. Otherwise, as shown by M. Consoli in this workshop [1], the strong evidence for triviality of $\lambda\phi^4$ theories leads to a non perturbative scheme without self interacting scalar sector and with a heavy Higgs of mass $M_H^2 = 8\pi^2 v^2$, $v$ being the vacuum expectation value of Higgs field.

In what follows we will assume that the fundamental theory of electroweak interactions is not precisely known, but that the possible symmetries are, i.e. gauge invariance and custodial symmetry.

We are interested in performing a spontaneous symmetry breaking avoiding physical scalar particles, i.e. by a non linear realisation. We will first explicitely show that the SM appears as a gauged $SU(2)_L \otimes SU(2)_R$ non linear $\sigma$-model [2]. Since this theory is non renormalisable it corresponds to an effective one. The use of effective Lagrangians dates from beggining of 60's with the introduction of the non linear $\sigma$-model as an effective theory for low energy strong interactions, exhibiting a chiral symmetry breakdown [3].

In what follows we will extensively use the fact that any non linear $\sigma$-model is gauge equivalent to theories with additional hidden local symmetry [4]. This has been shown in the context of pion interactions successfull to describe the $\rho$ vector resonances, which correspond to the gauge bosons associated to this additional hidden symmetry. Applying this mechanism to weak interactions we will build vector and axial vector resonances as gauge bosons associated to the hidden symmetry group of $SU(2)$ type. Under the assumption they are dynamical, we will get the $SU(2)$ BESS Lagrangian [5] (BESS standing for Breaking Electroweak Symmetry Strongly).

The existence of these new bosons called $\overrightarrow{V}$ will indirectly manifest at LEP through deviations from SM expectations [6]. For this purpose a low energy effective theory valid for heavy resonances will be derived.

The new particles are naturally coupled to fermions through mixing between $\overrightarrow{W}$ and $\overrightarrow{V}$, although a direct coupling is possible. Deviations from SM trilinear and quadrilinear gauge bosons couplings are expected.

In the vector dominance approximation the BESS model corresponds to technicolor model [7] with a single technidoublet. In order to reproduce the one family technicolor model [8] one has to extend the original $SU(2)_L \otimes SU(2)_R$ symmetry group to $SU(8)_L \otimes SU(8)_R$. This leads to the extended BESS model and to the existence of pseudogoldstone bosons [9].

These models are already constrained from the precision electroweak data as we will see. Exploration of the usefulness of very energetic linear $e^+e^-$ colliders and LHC to detect the new particle spectrum will also be reviewed.

## 2 Standard model revisited as a gauged non linear $\sigma$-model

The non fermionic part of the SM is described by the Lagrangian:

$$\begin{aligned}\mathcal{L}_{WS} = \tfrac{1}{2g'^2} Tr\left(B_{\mu\nu}B^{\mu\nu}\right) + \frac{1}{2g^2} Tr\left(F_{\mu\nu}(W)F^{\mu\nu}(W)\right) + D_\mu\phi^+ D^\mu\phi \\ -\lambda\left(\phi^+\phi + \frac{\mu^2}{2\lambda}\right)^2\end{aligned} \tag{1}$$

where $B_{\mu\nu}$ (resp $F_{\mu\nu}$) is the field strength of the $U(1)$ (resp $SU(2)$) field and $\phi = \binom{\phi^+}{\phi^0}$ the Higgs doublet.

$$\begin{aligned} B_{\mu\nu} &= \partial_\mu B_\nu - \partial_\nu B_\mu \quad \text{where} \quad B_\mu = ig'\tau_3 \frac{B'_\mu}{2} \\ W_{\mu\nu} &= \partial_\mu W_\nu - \partial_\nu W_\mu + [W_\mu, W_\nu] \quad \text{where} \quad W_\mu = ig\frac{\vec{\tau}}{2}\cdot \overrightarrow{W}'_\mu \end{aligned} \tag{2}$$

$D_\mu = \partial_\mu + W_\mu + B_\mu$ is the covariant derivative.

Let us assume $g = g' = 0$. The scalar part of (1) reads:

$$\mathcal{L}_S = \partial_\mu\phi^+\partial^\mu\phi - \lambda\left(\phi^+\phi + \frac{\mu^2}{\lambda}\right)^2 \tag{3}$$

exhibiting a global $SU(2)_L \otimes SU(2)_R$ global symmetry. Indeed introducing a matricial notation:

$$M = \sqrt{2}\left(i\tau_2\phi^\star,\ \phi\right) \tag{4}$$

the expression (3) becomes:

$$\mathcal{L}_S = \frac{1}{4} Tr\left(\partial_\mu M^+ \partial^\mu M\right) - \frac{\lambda}{4}\left(\frac{1}{2}Tr(M^+M) + \frac{\mu^2}{2\lambda}\right)^2 \tag{5}$$

invariant under

$$M \to M' = LMR^+ = e^{i\,\overrightarrow{\varepsilon_L}\cdot\overrightarrow{\tau}/2} M e^{-i\,\overrightarrow{\varepsilon_R}\cdot\overrightarrow{\tau}/2} \tag{6}$$

If one assume $\mu^2 < 0$ the potential has a minimum for

$$M^+M = -\frac{\mu^2}{\lambda} \equiv v^2 \tag{7}$$

leading to a spontaneous breaking of $SU(2)_L \otimes SU(2)_R$ into $SU(2)_D$.

The scalar part of the Lagrangian becomes:

$$\mathcal{L}_S = \frac{1}{4} Tr \partial_\mu M^+ \partial^\mu M - \frac{\lambda}{4}\left(\frac{1}{2} Tr M^+ M - v^2\right)^2 \tag{8}$$

Using

$$M = \sigma + i\overrightarrow{\pi} \cdot \overrightarrow{\tau} \tag{9}$$

the standard model is a linear $\sigma$ model

$$\mathcal{L}_S = \frac{1}{2}\partial_\mu \sigma \partial^\mu \sigma + \frac{1}{2}\partial_\mu \overrightarrow{\pi} \partial^\mu \overrightarrow{\pi} - \frac{\lambda}{4}\left(\sigma^2 + \overrightarrow{\pi}^2 - v^2\right)^2 \tag{10}$$

where $\sigma$ is the Higgs field.

Let us take now the limit $\lambda \to \infty$. The generating functionnal [10] :

$$\begin{aligned} G[\lambda] = \textstyle\int [d\overrightarrow{\pi}][d\sigma]\ & e^{i\int d^4x \left(\frac{1}{2}\partial_\mu \sigma \partial^\mu \sigma + \frac{1}{2}\partial_\mu \overrightarrow{\pi} \partial^\mu \overrightarrow{\pi}\right)} \\ & e^{-i\int d^4x \frac{\lambda}{4}\left(\sigma^2 + \overrightarrow{\pi}^2 - v^2\right)^2} \end{aligned} \tag{11}$$

becomes

$$G[\lambda] \underset{\lambda\to\infty}{=} \int [d\overrightarrow{\pi}][d\sigma]\ \delta\left(\sigma^2 + \overrightarrow{\pi}^2 - v^2\right) e^{i\int d^4x \left(\frac{1}{2}\partial_\mu \sigma \partial^\mu \sigma + \frac{1}{2}\partial_\mu \overrightarrow{\pi} \partial^\mu \overrightarrow{\pi}\right)} \tag{12}$$

leading to the non linear $\sigma$ model Lagrangian:

$$\mathcal{L}_S = \frac{1}{2}\partial_\mu \overrightarrow{\pi} \partial^\mu \overrightarrow{\pi} + \frac{1}{2}\frac{\left(\overrightarrow{\pi}\partial_\mu \overrightarrow{\pi}\right)\left(\overrightarrow{\pi}\partial^\mu \overrightarrow{\pi}\right)}{v^2 - \overrightarrow{\pi}^2} \tag{13}$$

which can be rewritten as:

$$\mathcal{L}_S = \frac{v^2}{4} Tr\left(\partial_\mu U \partial^\mu U^+\right) \tag{14}$$

where $U$ is an unitary matrix.

We introduce the Goldstone bosons $\varphi$ by parametrizing $U$ as:

$$U = e^{i\overrightarrow{\varphi}(x)\cdot\overrightarrow{\tau}/v} \tag{15}$$

Adding now the bosonic fields, the SM is a gauged non linear $\sigma$ model:

$$\mathcal{L}_B = \frac{v^2}{4} Tr\left(D_\mu U D^\mu U^+\right) + \frac{1}{2g'^2} Tr\left(B_{\mu\nu} B^{\mu\nu}\right) + \frac{1}{2g^2} Tr\left(F_{\mu\nu}(W) F^{\mu\nu}(W)\right) \tag{16}$$

Going into unitary gauge $U = 1$ one gets the mass terms for gauge bosons.

# 3 Hidden gauge symmetry

The basic idea of hidden gauge symmetry is that any non linear $\sigma$ model defined on coset space $G/H$ is gauge equivalent to $G \otimes H$ local where $H$ local is the hidden gauge symmetry group [11].

Let us write

$$U = LR^{+} \tag{17}$$

$L(x) \in SU(2)_L$ and $R(x) \in SU(2)_R$ are the Goldstone bosons. We define:

$$\rho(x) = \begin{pmatrix} L(x) & 0 \\ 0 & R(x) \end{pmatrix} \tag{18}$$

and introduce the Maurer Cartan form:

$$\omega_\mu(x) = \rho^{+}(x)\partial_\mu \rho(x) = \begin{pmatrix} L^{+}\partial_\mu L & 0 \\ 0 & R^{+}\partial_\mu R \end{pmatrix} \tag{19}$$

We split into:

$$\omega_\mu^{\prime\prime} = \frac{1}{2}\begin{pmatrix} L^{+}\partial_\mu L + R^{+}\partial_\mu R & 0 \\ 0 & L^{+}\partial_\mu L + R^{+}\partial_\mu R \end{pmatrix} \tag{20}$$

and

$$\omega_\mu^{\perp} = \frac{1}{2}\begin{pmatrix} L^{+}\partial_\mu L - R^{+}\partial_\mu R & 0 \\ 0 & -\left(L^{+}\partial_\mu L - R^{+}\partial_\mu R\right) \end{pmatrix} \tag{21}$$

Under local $SU(2)_V$ transformations $h(x)$ they behave as:

$$L^{+}\partial_\mu L \pm R^{+}\partial_\mu R \rightarrow h^{+}\left(L^{+}\partial_\mu L \pm R^{+}\partial_\mu R\right)h + \left(h^{+}\partial_\mu h \pm h^{+}\partial_\mu h\right) \tag{22}$$

i.e.,

$$\omega_\mu^{\perp} \rightarrow \widetilde{h}^{+}\omega_\mu^{\perp}\widetilde{h} \quad \text{and} \quad \omega_\mu^{\prime\prime} \rightarrow \widetilde{h}^{+}\omega_\mu^{\prime\prime}\widetilde{h} + \widetilde{h}^{+}\partial_\mu \widetilde{h} \tag{23}$$

with $\widetilde{h} = \begin{pmatrix} h & 0 \\ 0 & h \end{pmatrix}$

$\omega_\mu^{\prime\prime}$ transforms under $SU(2)_V$ as a gauge bosons $\mathcal{V}_\mu$:

$$\mathcal{V}_\mu \rightarrow h^{+}\mathcal{V}_\mu h + h^{+}\partial_\mu h \tag{24}$$

where $\mathcal{V}_\mu = i g_V \frac{\overrightarrow{\tau}}{2} \cdot V_\mu^{\prime}$.

The most general Lagrangian one gets from (19) is:

$$\mathcal{L} = -\frac{v^2}{2}\left[Tr\left(\omega_\mu^{\perp}\omega^{\perp\mu}\right) + \alpha\ Tr\left(\left(\omega_\mu^{\prime\prime} - \widetilde{\mathcal{V}}_\mu\right)\left(\omega^{\prime\prime\mu} - \widetilde{\mathcal{V}}^{\mu}\right)\right)\right] \tag{25}$$

where $\widetilde{\mathcal{V}}_\mu = \begin{pmatrix} \mathcal{V}_\mu & 0 \\ 0 & \mathcal{V}_\mu \end{pmatrix}$ and $\alpha$ is an arbitrary parameter.

The SM corresponds to the first term of (25) and $\widetilde{\mathcal{V}}_\mu = \omega''_\mu$ is an auxiliary field.

In order to enable these gauge bosons to show up as physical particles we have to add to the previous Lagrangian a kinetic term:

$$\mathcal{L}_{KIN} = \frac{1}{2g_V^2} Tr\left(F_{\mu\nu}(\mathcal{V}) F^{\mu\nu}(\mathcal{V})\right) \tag{26}$$

This term can be generated for $2D$ and $3D$ theories [12] and from quantum corrections for $4D$ theories [13]. To get the SM we have to perform a gauging of $SU(2)_L \otimes SU(2)_R$ into $SU(2)_L \otimes U(1)_Y$ and add kinetic terms for $W$ and $B$ fields (see eq. (1)). One gets [14] a Yang Mills Lagrangian whose gauge group is $SU(2)_L \otimes U(1)_Y \otimes SU(2)_V$ after replacement of the derivatives by the covariant ones:

$$D_\mu L = \left(\partial_\mu + W_\mu - \mathcal{V}_\mu\right) L$$

and

$$D_\mu R = \left(\partial_\mu + B_\mu - \mathcal{V}_\mu\right) R.$$

In the unitary gauge ($L = 1,\ R = 1$) one gets:

$$\begin{aligned}
\mathcal{L}_{BESS} = & -\frac{v^2}{4}\left[Tr(W-B)^2 + \alpha\ Tr(W+B-V)^2\right] \\
& + \frac{1}{2g^2} Tr\left(F_{\mu\nu}(W) F^{\mu\nu}(W)\right) + \frac{1}{2g'^2} Tr\left(B_{\mu\nu} B^{\mu\nu}\right) \\
& + \frac{1}{2g''^2} Tr\left(F_{\mu\nu}(V) F^{\mu\nu}(V)\right)
\end{aligned} \tag{27}$$

with $F_{\mu\nu}(V) = \partial_\mu V_\nu - \partial_\nu V_\mu + \frac{1}{2}[V_\mu,\ V_\nu]$ where $V_\mu = ig''\frac{\overrightarrow{\tau}}{2} \cdot \overrightarrow{V}'_\mu$ after the rescaling $g_V \equiv \frac{g''}{2}$. The first term within brackets is the usual mass term appearing the SM.

The fermionic part of the Lagrangian reads:

$$\begin{aligned}
\mathcal{L}_F = & i\overline{\psi}_L \gamma^\mu \left(\partial_\mu + W_\mu + i\left(Q - \frac{\tau_3}{2}\right) g' B'_\mu\right) \psi_L \\
& + i\overline{\psi}_R \gamma^\mu \left(\partial_\mu + iQg' B'_\mu\right) \psi_R \\
& + ib\overline{\psi}_L \gamma^\mu \left(\partial_\mu + \frac{V_\mu}{2} + i\left(Q - \frac{\tau_3}{2}\right) g' B'_\mu\right) \psi_L
\end{aligned} \tag{28}$$

where the last term corresponds to a direct coupling of the fermions to the fields $V$ and $Q$ is the electric charge.

The physical vector bosons $W^\pm, V^\pm, A, Z^0$ and $V^0$ are obtained after diagonalization of the charged and neutral sectors [14]. The mixing angles are of the order of $0\left(\frac{g}{g''}\right)$.

We will now study the low energy effects of the model.

# 4 Low energy effects of vector resonances

We will evaluate [15] the solution of the classical equations of motion in the limit $M_V \to \infty$.

From

$$\frac{\partial}{\partial V^a_\mu}\left(\mathcal{L}_{BESS} + \mathcal{L}_F\right) = 0 \tag{29}$$

we get:

$$g'' V^a_\mu = g W^a_\mu + g' B'_\mu \delta_{a_3} + \frac{b}{v^2 \alpha} \overline{\psi}_L \gamma^\mu \psi_L \tau_a \tag{30}$$

The last term in eq. (30) will be neglected since $b$ is small.

The physical effect of $V$ to low energy is present in the $V$ kinetic term and the interacting fermionic Lagrangian where one has to replace $V_\mu$ by expression (30).

Neglecting for the moment the trilinear and quadrilinear couplings one gets:

$$\begin{aligned}\mathcal{L}_{eff} = &-\frac{1}{4}(1+Z_\gamma)A_{\mu\nu}A^{\mu\nu} - \frac{1}{2}(1+Z_W)W^+_{\mu\nu}W^{\mu\nu -}\\ &-\frac{1}{4}(1+Z_Z)Z_{\mu\nu}Z^{\mu\nu} + \frac{1}{2}Z_{Z_\gamma}A_{\mu\nu}Z^{\mu\nu}\\ &-M_W^2 W^+_\mu W^{\mu -} - \frac{1}{2}M_Z^2 Z_\mu Z^\mu\end{aligned} \tag{31}$$

with:

$$\begin{aligned}Z_\gamma &= \left(\frac{g}{g''}\right)^2 4\sin^2\theta_W\\ Z_W &= \left(\frac{g}{g''}\right)^2\\ Z_Z &= \left(\frac{g}{g''}\right)^2 \frac{\cos^2 2\theta_W}{\cos^2\theta_W}\\ Z_{Z_\gamma} &= -\left(\frac{g}{g''}\right)^2 2\,\mathrm{tg}\,\theta_W \cos 2\theta_W\end{aligned} \tag{32}$$

The corrections to the SM produce a wave function renormalization of the fields $A_\mu, Z_\mu$ and $W^\pm_\mu$ and a mixing term $Z_{Z\gamma}$ to be cancelled.

Defining new fields indiced by the superscript $R$:

$$\begin{aligned}A_\mu &\to \left(1 - \frac{Z_\gamma}{2}\right)A^R_\mu + Z_{Z_\gamma} Z^R_\mu\\ W_\mu &\to \left(1 - \frac{Z_W}{2}\right)W^R_\mu\\ Z_\mu &\to \left(1 - \frac{Z_Z}{2}\right)Z^R_\mu\end{aligned} \tag{33}$$

the only deviations are present in the mass terms of eq. (31) i.e.

$$M_W^2 \to M_W^2(1-Z_W) \quad \text{and} \quad M_Z^2 \to M_Z^2(1-Z_Z). \tag{34}$$

The couplings of the gauge bosons to fermions will also be affected:

$$\mathcal{L}_{em} = -e\left(1-\frac{Z_\gamma}{2}\right)\overline{\psi}\gamma^\mu Q\psi A_\mu^R \tag{35}$$

$$\mathcal{L}_{charged} = -\frac{e}{\sqrt{2}\sin\theta_W}W_\mu^R\overline{\psi}_u\gamma^\mu\frac{1-\gamma_5}{2}\psi_d\left(1-\frac{b}{2}-\frac{Z_W}{2}\right)+h.c. \tag{36}$$

$$\begin{aligned}\mathcal{L}_{neutral} = &-\frac{e}{\sin\theta_W\cos\theta_W}\left(1-\frac{b}{2}-\frac{Z_Z}{2}\right)\overline{\psi}\gamma^\mu\left[T_{3L}\frac{1-\gamma_5}{2}\right.\\ &\left.-Q\sin^2\theta_W\left(1+\frac{b}{2}-\operatorname{cotg}\theta_W Z_{Z_\gamma}\right)\right]\psi Z_\mu^R\end{aligned} \tag{37}$$

The input parameters that are used for LEP Physics are:

— the electric charge $e^{ph}$

— the Fermi constant $G_F^{ph}$

— the mass of the $Z$ boson $M_Z^{ph}$

They are identified as follows:

$$e^{ph} = e\left(1-\frac{Z_\gamma}{2}\right) \tag{38}$$

$$M_Z^{2ph} = M_Z^2\,(1-Z_Z) \tag{39}$$

$$\frac{G_F}{\sqrt{2}} = \frac{e^{2ph}\,(1-b-Z_Z+Z_\gamma)}{8\sin^2\theta_W\cos^2\theta_W M_Z^{2ph}} \tag{40}$$

Since

$$\frac{G_F}{\sqrt{2}} = \frac{4\pi\alpha(M_Z)}{8\sin^2\theta_W^{SM}\cos^2\theta_W^{SM}M_Z^{2ph}} \tag{41}$$

we can connect $\sin\theta_W$ to the tree level SM value $\sin\theta_W^{SM}$.

The strength of corrections due to $V$ to the ratio:

$$\frac{M_W^2}{M_Z^2} = \frac{1}{2}+\sqrt{\frac{1}{4}-\frac{\pi\alpha(M_Z)}{\sqrt{2}G_F M_Z^2(1-\Delta r_W)}} \tag{42}$$

is given by

$$\Delta r_W = 2\left(\frac{g}{g''}\right)^2 - b \tag{43}$$

The neutral Lagrangian (37) can be rewritten as:

$$\mathcal{L}_{neutral} = -\frac{e^{ph} Z_\mu}{\sin\theta_W^{SM}\cos\theta_W^{SM}}\overline{\psi}\gamma^\mu\left(g_V + \gamma^5 g_A\right)\psi \tag{44}$$

with

$$\begin{aligned} g_V &= \frac{T_{3_L}}{2} - Q\sin^2\overline{\theta}_W \\ g_A &= -\frac{T_{3_L}}{2} \\ \sin^2\overline{\theta}_W &= (1+\Delta K)\sin^2\theta_W^{SM} \end{aligned} \tag{45}$$

where

$$\Delta K = \frac{1}{\cos 2\theta_W^{SM}}\left(\left(\frac{g}{g''}\right)^2 - \frac{b}{2}\right) \tag{46}$$

We get also

$$\Delta\rho = 0 \tag{47}$$

LEP 200 will directly test the non abelian gauge structure through the trilinear and quadrilinear vertices among gauge bosons.

The anomalous vertices read:

$$\begin{aligned} \mathcal{L}_{trilinear} &= \frac{ie^{ph}}{\sin^2\theta_W^{SM}\cos^2\theta_W^{SM}}\left(b\cos^2\theta_W^{SM} - \left(\frac{g}{g''}\right)^2\right) \\ &\quad \left(Z_{\mu\nu}W^{-\mu}W^{+\nu} + Z^\mu W^{-\nu}W^+_{\mu\nu} + Z^\nu W^{+\mu}W^-_{\mu\nu}\right) \end{aligned} \tag{48}$$

and

$$\begin{aligned} \mathcal{L}_{quadri} &= (2g_{\mu\rho}g_{\nu\sigma} - g_{\mu\nu}g_{\rho\sigma} - g_{\mu\sigma}g_{\rho\nu})\, W^{+\mu}W^{-\rho} \\ &\quad \left[-e^{2ph}\,\mathrm{cotg}\,\theta_W^{SM}\gamma_1 A^\nu Z^\sigma + \frac{1}{2}\frac{e^{2ph}}{\sin^2\theta_W^{SM}}\gamma_2 W^{+\nu}W^{-\sigma}\right. \\ &\quad \left. -\frac{1}{2}e^{2ph}\,\mathrm{cotg}^2\,\theta_W^{SM}\gamma_3 Z^\nu Z^\sigma\right] \end{aligned} \tag{49}$$

with

$$\gamma_1 = \frac{1}{2\cos^2\theta_W^{SM}\cos 2\theta_W^{SM}}\left(b\cos^2\theta_W^{SM} - \left(\frac{g}{g''}\right)^2\right)$$

$$
\begin{aligned}
\gamma_2 &= \frac{\cos^2\theta_W^{SM}}{\cos 2\theta_W^{SM}} b + \left(\frac{g}{g''}\right)^2 \left(\frac{1}{4} - \frac{1}{\cos 2\theta_W^{SM}}\right) \\
\gamma_3 &= \frac{b}{\cos 2\theta_W^{SM}} - \left(\frac{g}{g''}\right)^2 \frac{1 + 2\cos^2\theta_W^{SM}}{4\cos^4\theta_W^{SM}\cos 2\theta_W^{SM}}
\end{aligned}
$$

We are now ready to derive the bounds on the BESS parameter space coming from precise LEP measurements and study the potential discovery at future colliders.

# 5 LEP constraints

The analysis of LEP data, concerning the total width, the hadronic width, the leptonic width, the leptonic and $b$ forward-backward asymmetries, the $\tau$-polarization, the cesium atomic parity violation and the ratio $M_W/M_Z$, uses available full one loop radiative correction programs [16]. This brings in a dependance on $\alpha_S$, $m_{top}$ and a cut off $\Lambda$ which corresponds to the Higgs mass in the Standard Model. The quantities $\Delta K$, $\Delta\rho$, $\Delta r_W$ are directly connected to observable quantities. We will reexpress them in terms of the parameters $\varepsilon_i$ [17] :

$$
\begin{aligned}
\varepsilon_1 &= \Delta\rho \\
\varepsilon_2 &= \cos^2\theta_W \Delta\rho + \frac{\sin^2\theta_W}{\cos 2\theta_W}\Delta r_W - 2\sin^2\theta_W \Delta K \\
\varepsilon_3 &= \cos^2\theta_W \Delta\rho + \cos 2\theta_W \Delta K
\end{aligned} \tag{50}
$$

The BESS contribution reads:

$$
\begin{aligned}
\varepsilon_1 &= \varepsilon_2 = 0 \\
\varepsilon_3 &= \left(\frac{g}{g''}\right)^2 - \frac{b}{2}
\end{aligned} \tag{51}
$$

This shows explicitely that through LEP data we are only sensitive to one combination of BESS parameters i.e. $\varepsilon_3$. The allowed region at 90% CL in the $\left(\left(\frac{g}{g''}\right)^2,\ \frac{b}{2}\right)$ plane is shown in fig. 1 for three top mass values.

The chosen experimental value [18]

$$
\varepsilon_3^{\exp} = (3.4 \pm 1.8)10^{-3} \tag{52}
$$

corresponds to La Thuile et Moriond data.

The two standard deviation from Standard Model expectation for $b$ partial width can be expressed in terms of $\varepsilon_b$ parameter [19].

Assuming a non zero direct coupling only for the heaviest generation (as expected from one loop BESS radiative corrections proportional to $m_f$) we get:

$$\varepsilon_b = -\frac{b}{2} \tag{53}$$

After adding the SM expectation for $m_t = 170$ GeV we get at 90% CL

$$-3.0\ 10^{-2} \leq b \leq -3.4\ 10^{-3}. \tag{54}$$

# 6 Discovery potential at future colliders

Provided the center of mass energy is higher than $2M_W$ we are directly sensitive to trilinear and quadrilinear gauge bosons couplings. LEP 200 energy is too small to be really sensitive to the expected BESS model deviations. Fortunately the planned linear $e^+e^-$ colliders will put severe constraints especially through the reaction $e^+e^- \to W^+W^-$ which deviates from SM values due to $V^0$ exchange [20]. The best constraints are obtained from longitudinally polarized $W_L^+W_L^-$ final state as shown in fig. 2. Needless to remind that if a $V^0$ resonance exists below the center of mass energy it will show up in $e^+e^- \to f\overline{f},\ W^+W^-$.

Hadronic colliders as especially the LHC are well suited to discover the charged resonances $V^\pm$ [21].

Two subprocesses contribute: the quark antiquark annihilation and the $\gamma W$ or $ZW$ fusion. The appropriate final state is $WZ$ followed by leptonic decays which is not affected by the top background since it can be reduced from $Z$ mass reconstruction of lepton pair and use of isolation criteria for leptons [22].

The charged resonances show up as broad resonances around the $V^\pm$ mass in $WZ$ invariant mass spectrum or a broad Jacobian peak around $\frac{M_V}{2}$ in the $Z$ transverse momentum spectrum.

At LHC, assuming an integrated luminosity of $10^5 pb^{-1}$/year, one can reach discovery limits up to 2 TeV. Nevertheless this limit cannot be reached for the full left over $\left(\frac{g}{g''},\ b\right)$ parameter space. Fig. 3 shows the favorable choice ($g'' = 20,\ b = 0.0\ 16$) for $M_V = 1.5$ TeV whereas as can be inferred from fig. 4 for $g'' = 20$ and $b = 2.0\ 10^{-3}$ the signal exhibits no singular behaviour from background.

# 7 Extended BESS model

One important specialization of BESS model is to technicolor theories since for particular values of the parameters it would correspond to a technicolor model involving a single technidoublet ($N_d = 1$). If a non zero direct coupling of $V$ to fermions exists it corresponds to an extended technicolor [23] with:

$$b = -2\left(\frac{v}{\Lambda_{ETC}}\right)^2 \tag{55}$$

where $\Lambda_{ETC}$ is the scale for extended technicolor.

Previous analyses have shown that a conventional QCD scaled technicolor is excluded for $N_{TC}N_d \leq 12$ at 90% CL. Therefore we will extend the original non linear $SU(2)_L \otimes SU(2)_R$ $\sigma$-model to a $SU(8)_L \otimes SU(8)_R$ broken to $SU(8)_V$ in order to incorporate the one family technicolor model ($N_d = 4$). The model [9] will contain not only vector resonances but also axial vector resonances and pseudogoldstone bosons.

The construction starts from a gauged non linear $\sigma$-model of gauge group $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$:

$$\mathcal{L} = \frac{v^2}{16} Tr\left[(D_\mu U) D^\mu U^+\right] \tag{56}$$

with

$$U = \exp\left(\frac{2i\pi^A \cdot T^A}{v}\right)$$

The $\pi^A$ are the Goldstone bosons whereas the $T^A$ are the $SU(8)$ generators:

$$T^A = \left(T^a,\ \widetilde{T}^a,\ T^D,\ T_8^\alpha,\ T_8^{a\alpha},\ T_3^{\mu i},\ \widetilde{T}_3^{\mu i}\right)$$

where $A = 1, \cdots 63$, $a = 1,2,3$ is a $SU(2)$ triplet index, $\alpha = 1, \cdots 8$ corresponds to $SU(3)_C$ octet indices, $i = 1,2,3$ to $SU(3)_C$ triplet indices and $\mu = (0, a)$. $T^D$ is a singlet under $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$.

The SM gauge fields can be written as:

$$\begin{aligned}
\mathcal{A}_\mu &= 2igW_\mu^a T^a + i\sqrt{2} g_S G_\mu^\alpha T_8^\alpha + 2i\frac{1}{\sqrt{3}} g' B_\mu T^D \\
\mathcal{B}_\mu &= 2ig' B_\mu \left(T^3 + \frac{1}{\sqrt{3}} T^D\right) + i\sqrt{2} g_S G_\mu^\alpha T_8^\alpha
\end{aligned} \tag{57}$$

and the new resonances:

$$\begin{aligned}
V_\mu &= ig'' V_\mu^A T^A \\
A_\mu &= ig'' A_\mu^A T^A
\end{aligned}$$

The most general Lagrangian for gauge bosons reads:

$$\begin{aligned}
\mathcal{L} = -\frac{v^2}{16} &\left(a\ Tr(\mathcal{A} - \mathcal{B})^2 + b\ Tr(\mathcal{A} + \mathcal{B} - 2V)^2\right. \\
&\left. + cTr(\mathcal{A} - \mathcal{B} + 2A)^2 + d\ Tr(2A)^2\right) \\
&+ \text{kinetic terms}
\end{aligned} \tag{58}$$

The $A$ mass is a new parameter and we will assume no direct coupling to fermions.

The pseudo-Goldstone mass spectrum has been derived [24] from the one loop effective potential, which includes, besides the ordinary gauge interactions, the Yukawa couplings. The masses are proportional to the ultraviolet cut-off $\Lambda$ and depending on the heaviest fermions masses.

The restrictions from LEP precision measurements are obtained from the $\varepsilon_i$ parameters. The parameters $\varepsilon_1$ and $\varepsilon_2$, which are isospin violating, are sensitive to the pseudogoldstone mass spectrum. The effect on the parameter $\varepsilon_1$ is to weaken the upper bound on the top mass. The contribution to $\varepsilon_3$ depends on the cut-off $\Lambda$ and on the masses of the pseudogoldstone, vector and axial vector bosons. Except for small pseudogoldstone boson masses, it is negative.

Concerning hadronic colliders the discovery limit of charged resonances through $WZ$ final state is lowered compared to $SU(2)$ case as soon as decay into pseudogoldstones bosons is allowed [25]. If these particles can be copiously produced at LHC they have to suffer a huge background since $P^0$ (resp $P^{\pm}$) decay into $t\bar{t}$ or $b\bar{b}$ (resp $t\bar{b}$ and $\bar{t}b$). These backgrounds have been studied for charged Higgs boson discovery from $tb$ decays at LHC using SDC detector. It has been shown that reasonably efficient and pure $b$ tagging is mandatory. Our case deserves a careful study along the previous procedure to be conclusive.

The most promising pseudogoldstone pair production mechanism at $e^+e^-$ linear colliders is the resonant one through a $V^0$, which is accompanied by an enhancement of $W$ pair production. The final state to be considered is $t\bar{b}\bar{t}b$, already considered for charged Higgs production.

# 8 Conclusion

We have studied the possibility of a strong interacting sector being at the origin of the electroweak symmetry breaking, avoiding elementary scalars. In absence of a specific definite theory of the strong electroweak sector the BESS model provides for a rather general frame based on custodial symmetry and gauge invariance.

A characteristic feature of the model is the occurence of spin one resonances in the $TeV$ range and pseudogoldstone bosons in its extended version. The idea of a strong electroweak sector is quantitatively testable at LEP thanks to the recent precision measurements: QCD scaled technicolor has already been excluded.

Linear $e^+e^-$ colliders are sensitive to the neutral vector resonance $V^0$ whereas energetic hadronic colliders are well suited to discover the charged ones. The characteristic feature of the extended BESS model is the existence of pseudogoldstone bosons whose detection needs a careful evaluation of backgrounds which remain to be done.

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